

# Realistic GUT with Gauge Mediated Supersymmetry Breaking

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## Abstract

We present an example of the gauge mediated SUSY breaking flipped  $SU(5)$  model. The messengers of the SUSY breaking are either only colour triplets which belong to the minimal content of the scalar supermultiplets or together with triplets as a messengers emerge the ordinary Higgs doublets. In both cases the model predicts light gauginos in respect of the squarks and sleptons, which could be tested in the nearest LEP experiments.

In both cases "all order" solution of the doublet-triplet splitting problem is obtained, the  $\mu$ -term of the order of 100 GeV is generated and the left handed neutrino masses are suppressed.

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The supersymmetric theories suggest the elegant possibilities for solution of the gauge hierarchy problem. Non-renormalization theorems [1] in SUSY theories imply that certain ratios of coupling constants are non-renormalized in exact SUSY limit. This nice feature and also the successful prediction of the numerical value of  $\sin^2 \theta_W$  [2] supported the idea of the SUSY Grand Unified Theories (GUT).

The most interesting question is the origin of the SUSY breaking. It is usually assumed that SUSY is broken in a "hidden" sector and by some interactions transmitted in the visible sector. The most famous scenarios are the supergravity theories [3], in which the SUSY breaking in the visible sector transmitted by the gravity. In this case the soft SUSY breaking (SSB) terms are presented at the energies which correspond to the Planck scale -  $M_P$  and even they have the universal form, they will renormalized between the  $M_P$  and  $M_{GUT}$ . At  $M_{GUT}$  one has to integrate out the heavy particles and evolve again the SSB parameters from  $M_{GUT}$  to  $m_W$  scale with the RGEs of the MSSM. These processes violate the universalities (see [4] and references there) and lead to the flavour changing neutral currents (FCNC).

To another class of the SUSY breaking scenarios belong the gauge mediated supersymmetry breaking (GMSB) models [5, 6], in which the supersymmetry breaking is transmitted by the gauge interactions. Because the fact, that this models do not suffer from the FCNC problem the interest in models of this type was renewed recently [7, 8].

In this paper we present an example of the SUSY GUT in which the SUSY breaking occurs in the sector of scalar superfields which are used for the GUT symmetry breaking.

The main contribution to the soft masses to the squarks and sleptons comes from the nonzero  $D$ -term, which is just of the order of SUSY scale; while gauginos gain masses through the  $SU(3)_C \times SU(2)_W \times U(1)_Y$  gauge interactions.

As a realistic model we consider the flipped  $SU(5)$  theory which provides the natural solution of the doublet-triplet (DT) splitting problem through the missing partner mechanism. Crucial role in the SUSY breaking is played by the anomalous  $U(1)_A$  symmetry.

The model predicts gauginos with masses in the range of 1 GeV, while the soft masses of the scalar particles are in the region of  $10^2 - 10^3$  GeV.

By the special implementation of the  $U(1)_A$  charges of some superfields the model suggests two different sets of the messenger superfields. In first case in the role of the messenger superfields emerge the colour triplets and standard electroweak Higgs doublets, which could give the nonuniversal contributions to the masses of the squarks and sleptons; However in the case considered in the present paper these contributions are strongly suppressed. In this case the masses of all gauginos are generated through the one loop diagrams. While in the second case in the role of the messengers we naturally have only color triplets and no mass term generated for wino. In both cases the desirable  $\mu$ -term is generated, left handed neutrinos are naturally light and proton decay through the  $d = 5$  operators is strongly suppressed.

Let us note, that some examples in which the standard Higgs doublets emerge as a messengers of the SUSY breaking also was considered in the recent work [9] .

The flipped version of the  $SU(5)$  model provides the solution of the DT splitting problem through the missing partner mechanism by the most economical way [10, 11] .

The gauge group is  $SU(5) \times U(1)$  and the matter superfields transform under this group as:  $(10_1 + \bar{5}_{-3} + 1_5)_i$  ( $i$  is a family index ) in which the ordinary quark and lepton superfields are compressed as:

$$\begin{aligned} 10_1 &= (q, d^c, \nu_R)_1 , \\ \bar{5}_{-3} &= (u^c, l)_{-3} , \\ 1_5 &= e_5^c . \end{aligned} \tag{1}$$

The  $10_1$  contains  $\nu_R$  additional state which is singlet under the  $G_{321} \equiv SU(3)_C \times SU(2)_W \times U(1)_Y$  group. The Higgs sector consists to the following superfields:

$$\begin{aligned} H \sim 10_1 &= (Q, D^c, N)_1 , & \bar{H} \sim \bar{10}_{-1} &= (\bar{Q}, \bar{D}^c, \bar{N})_{-1} , \\ \phi \sim 5_{-2} &= (T, h_d)_{-2} , & \bar{\phi} \sim \bar{5}_2 &= (\bar{T}, h_u)_2 . \end{aligned} \tag{2}$$

The  $h_u$  and  $h_d$  fields are generate the masses of up and down quarks and leptons:

$$W_Y^0 = 10 \cdot \bar{5} \cdot \bar{\phi} + 10 \cdot 10 \cdot \phi + 1 \cdot \bar{5} \cdot \phi . \tag{3}$$

the first term generates the masses of up type quarks, while the second and third - masses of down quarks and leptons respectively.

The  $H + \bar{H}$  pair is used for the GUT symmetry breaking. If  $N + \bar{N}$  from the set  $H + \bar{H}$  develop VEVs of order  $M_X \simeq 10^{16}$  GeV, then  $SU(5) \times U(1)$  directly is broken to  $G_{321}$  and  $Q(3, 2) + \bar{Q}(\bar{3}, 2)$  from  $H + \bar{H}$  are eaten Goldstone modes. The couplings between  $H$  ( $\bar{H}$ ) and  $\phi$  ( $\bar{\phi}$ ) superfields are described by the superpotential:

$$W_1^0 = \lambda_1 H H \phi + \lambda_2 \bar{H} \bar{H} \bar{\phi} . \tag{4}$$

Substituting the VEVs of the  $N$  and  $\bar{N}$  fields the mass terms for the triplet components will get the form:

$$W_m = \lambda_1 \langle N \rangle D^c T + \lambda_2 \langle \bar{N} \rangle \bar{D}^c \bar{T} . \tag{5}$$

So, after the GUT symmetry breaking the triplet states decouple. While  $H$  and  $\bar{H}$  do not contain doublet fragments,  $h_u$  and  $h_d$  remain naturally light.

Suppose, by some mechanism (which will be presented below)  $N$  and  $\bar{N}$  have F-terms with nonzero VEVs which magnitudes are of the order  $\sim m M_X$  ( $m$  is mass scale up to

which the SUSY is switched on). This will cause the shift between the masses of the scalar and fermionic components from the triplet superfields by the value  $\sim \sqrt{m\bar{M}_X}$ . While  $m\bar{M}_X \ll \langle N \rangle^2, \langle \bar{N} \rangle^2$ , masses of the scalar components will not be changed (see (5)). So, we will not have the light triplet states in the spectra and the successful unification of the three gauge coupling constants will not be altered. While the colour-triplet fragments transform nontrivially under  $SU(3)_C$  and  $U(1)_Y$  gauge groups, the SUSY breaking will be transferred from the  $W_1^0$  sector by the gauge interactions. The gauginos get masses through one loop diagrams. For a general set of messengers their masses are given by the formula [7]:

$$M_a = \frac{\alpha_a}{4\pi} \frac{F_i}{M} n_a(i) , \quad (6)$$

where the  $M$  is the mass of the corresponding messenger, and  $F_i$  - appropriate  $F$ -term.  $n_a(i)$  is Dynkin index and for fundamental representation of  $SU(N)$  equals to 1 and for  $U(1)_Y$ ,  $n_1 = \frac{6}{5}Y^2$  ( $Y = Q_{em} - T_3$ ). Index  $a$  in (6) corresponds to the gauge group and is 1, 2 and 3 for  $U(1)_Y$ ,  $SU(2)_W$  and  $SU(3)_C$  respectively. The (6) is written for  $F_i \ll M^2$  case.

In our case only gluinos and  $U(1)_Y$  gauginos get masses and winos remain massless in the one loop level, since there are no doublets among the messengers.

In this situation one state of chargino is lighter than  $W$  boson <sup>2</sup>. According to refs. [12] - [14] this case is not excluded and requires the low  $\tan\beta$  regime.

The squarks and sleptons get masses through two loop diagrams [15] and are given by the following formula:

$$\tilde{m}^2 = 2 \left( \frac{F}{M} \right)^2 \sum \left( \frac{\alpha_a}{4\pi} \right)^2 C_a n_a , \quad (7)$$

where  $C_3 = 4/3$ ,  $C_2 = 3/4$  and  $C_1 = 3/5Y^2$ .

Let us now describe how the SUSY breaking occurs in our model. For SUSY breaking we introduce the anomalous  $U(1)_A$  gauge symmetry. As a source of the SUSY breaking  $U(1)_A$  symmetry was used in the recent works [16]. The properties of the anomalous  $U(1)_A$  symmetry also were used for explaining the problem of gauge hierarchy [17, 18] as well as for the understanding of the pattern of fermion masses and mixing [18, 19].

It is well known, that anomalous extra  $U(1)$  factors appear in effective field theories from strings. The cancellation of its anomalies occurs by the Green-Schwarz mechanism

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<sup>2</sup>This happens if the  $\mu$  term exists for the doublet components. If  $\mu$ -term is zero we will have two light states in the theory. However as will be shown later in our model the  $\mu$  term with the desirable magnitude can be generated.

[20] . Because of the anomaly the Fayet-Iliopoulos term is always generated [21] and in string theories it equals to [22]

$$\xi = \frac{g_A^2 M_P^2}{192\pi^2} \text{Tr} Q . \quad (8)$$

So,  $D_A$  term will have the form:

$$\frac{g_A^2}{8} D_A^2 = \frac{g_A^2}{8} \left( \sum q_i |\varphi_i|^2 + \xi \right)^2 . \quad (9)$$

In our model, which gauge group is  $G = SU(5) \times U(1) \times U(1)_A$ , the superfields have the following prescription of the  $U(1)_A$  charges:  $q_H = q_{\bar{H}} = -1$ ,  $q_\phi = q_{\bar{\phi}} = 2$ . This choice of the charges will not change the form of the  $W_1^0$  (see (4)). We also introduce the singlet superfields  $X$ ,  $Y$  and  $\bar{Y}$  with  $U(1)_A$  charges:  $q_X = 2$  and  $q_Y = q_{\bar{Y}} = 0$ . The scalar superpotential

$$W_0 = \lambda \frac{\bar{Y} Y}{M_P^2} X \bar{H} H \quad (10)$$

is the most general under  $G \times \mathcal{R}$  symmetry, where  $\mathcal{R}$  symmetry acts on superfields  $\varphi_i \rightarrow e^{iR_{\varphi_i}} \varphi_i$  in such a way that  $W \rightarrow e^{i\alpha} W$ . So, the  $\mathcal{R}$  ‘charges’ of the superpotential and superfields are arranged as follows:

$$\begin{aligned} R_W = \alpha , \quad R_{\bar{Y}} = \alpha - 2R_{\bar{H}} , \quad R_\phi = \alpha - 2R_H , \\ R_Y + R_{\bar{Y}} = \alpha - R_H - R_{\bar{H}} - R_X , \end{aligned} \quad (11)$$

so,  $W_1^0 + W_0$  (see (4) and (10)) can be the most general without fixing  $R$  numbers of the all superfields. The potential builded from  $F$  and  $D$ -terms will have the form:

$$V = \sum |F_i|^2 + \tilde{g}^2 \left( |H|^2 - |\bar{H}|^2 \right)^2 + \frac{g_A^2}{8} \left( \xi - |H|^2 - |\bar{H}|^2 + 2|X|^2 \right)^2 , \quad (12)$$

where  $\tilde{g}^2 = \frac{3}{10}g^2 + \frac{1}{8}g_1^2$  ( $g$  and  $g_1$  are the  $SU(5)$  and  $U(1)$  coupling constants respectively). Supposing that  $\xi > 0$ , one can easily see that there exists the SUSY conserving minima:

$$\begin{aligned} \bar{Y} Y = 0 , \quad |H| = |\bar{H}| , \\ |H|^2 + |\bar{H}|^2 - 2|X|^2 = \xi ; \end{aligned} \quad (13)$$

So, for the scalar superpotential given by (10) SUSY remains unbroken.

Let us imply the proposal of ref. [16] and suppose that the  $\bar{Y}$  and  $Y$  superfields transform nontrivially under the some gauge group which interaction becomes strong below some  $\Lambda$  scale . The simplest case is the  $SU(2)$  gauge group under which  $Y$  and

$\bar{Y}$  are the pair of doublet-antidoublet. Non-perturbative superpotential induced by the instanton effect have the form [24] :

$$W_{\text{inst}} = \frac{\Lambda^5}{\bar{Y}Y} \quad (14)$$

and whole scalar superpotential will be <sup>3</sup> :

$$W = \lambda \frac{\bar{Y}Y}{M_P^2} X \bar{H} H + \frac{\Lambda^5}{\bar{Y}Y} . \quad (15)$$

The  $F$  and  $D$ -terms will have the forms:

$$\begin{aligned} F_H &= \lambda \frac{\bar{Y}Y}{M_P^2} X \bar{H} , & F_{\bar{H}} &= \lambda \frac{\bar{Y}Y}{M_P^2} X H , & F_X &= \lambda \frac{\bar{Y}Y}{M_P^2} \bar{H} H , \\ F_Y &= \lambda \frac{\bar{Y}}{M_P^2} X \bar{H} H - \frac{\Lambda^5}{\bar{Y}Y^2} , & F_{\bar{Y}} &= \lambda \frac{Y}{M_P^2} X \bar{H} H - \frac{\Lambda^5}{\bar{Y}^2 Y} \\ D &= |H|^2 - |\bar{H}|^2 , & D_A &= \xi - |H|^2 - |\bar{H}|^2 + 2|X|^2 . \end{aligned} \quad (16)$$

It is easy to see that SUSY is broken because there is no solution with vanishing  $F$  and  $D$  terms.

Minimizing the potential, builded from the  $F$  and  $D$  terms (see (16)), we can find that minimum can be obtained for the solutions:

$$\begin{aligned} H^2 &= \bar{H}^2 = \frac{3}{5}\xi , & X^2 &= \frac{1}{10}\xi + \frac{5m^2}{g_A^2} , \\ \bar{Y}^4 &= Y^4 = \frac{25}{3}M_P^4 \frac{m^2}{\xi} \left( 1 + \frac{125}{6\sqrt{3}\lambda} \frac{mM_P^2}{\xi\sqrt{\xi}} \right) . \end{aligned} \quad (17)$$

where

$$m^2 = \frac{2}{\sqrt{10}} \frac{\lambda \Lambda^5}{M_P^2 \sqrt{\xi}} \quad (18)$$

and for  $\Lambda \sim 10^{11} - 10^{12}$  GeV,  $\sqrt{\xi} \sim 10^{16}$  GeV and  $M_P \sim 10^{18}$  GeV we obtain  $m \sim 100$  GeV-10 TeV. For (17) solutions taking into account (16), (18)

$$F_X \sim F_H = F_{\bar{H}} = m\sqrt{\xi} , \quad F_Y = F_{\bar{Y}} \sim mM_P \left( \frac{m}{\sqrt{\xi}} \right)^{1/2} ,$$

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<sup>3</sup>Non-perturbative term can violate the R symmetry if the R symmetry is an anomalous.

$$D = 0, \quad D_A = \frac{10m^2}{g_A^2}. \quad (19)$$

As we see the SUSY in broken and nonzero  $F_H, F_{\bar{H}}$ -terms are the middle geometrical between the  $M_X$  (GUT scale) and  $\sim m_W$  scales ; Also the nonvanishing  $D_A$ -term with magnitude  $m^2$  is generated and the main contribution in the masses of the scalar components of the ordinary superfields comes from this term <sup>4</sup> and equals to:

$$\tilde{m}_{\varphi_i}^2 = \frac{5}{2}m^2 q_i, \quad (20)$$

where  $q_i$  is the anomalous  $U(1)_A$  charge of the appropriate  $\varphi_i$  superfield. The non-universal contributions to the squark masses through the supergravity corrections  $\sim m^2 \epsilon_X^2$  (where  $\epsilon_X = M_X/M_P$ ) for  $\epsilon_X \sim 10^{-2}$  will be negligible.

The upper bound of the soft masses (which are also proportional to the  $m^2$ ) of the electroweak Higgs doublets could be obtained from the requirement of the electroweak symmetry breaking and related to the mass of the top quark. Namely for  $m_t = 175 - 180$  GeV upper bound on  $m_{H_u}^2$  is  $\sim (350 \text{ GeV})^2$  [25]. For this order of  $m^2$  the mass of the gluino and also 'Majorana' masses of wino and zino are of the order of 1 GeV. The recent analyses of percentage exclusion of such a light gluino was presented in [27], where results of [13] were performed. Existence of light (or massless) wino leads to the one state of chargino lighter than  $W$  boson. According refs. [12, 14] this case did not excluded and requires the low  $\tan \beta$  regime. As far as the light bino concerned its phenomenological implications were described in refs. [12] .

The GMSB example with light gauginos was presented in ref. [26] and with light wino in ref. [14]; While in the recent [28] works the models with a gluino as a lightest SUSY particle (LSP) were considered.

While the gaugino masses are generated through the nonzero F-terms (of the  $N + \bar{N}$  components of the  $H + \bar{H}$  pair) in the one loop level, there magnitude will be  $\alpha_a/(4\pi)m$ . So, the model predicts the gauginos with low soft masses in respect to the soft masses of the squarks and sleptons.

From (20) we see that the matter superfields must have the positive  $U(1)_A$  charges. This can easily obtained if  $W_1^0$  (see (4)) will be rewritten to the form:

$$W_1 = \lambda_1 H H \phi + \lambda_2 \frac{Z}{M_P} \bar{H} \bar{H} \bar{\phi} \quad (21)$$

and the Yukawa superpotential for ordinary quarks and leptons will have the form:

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<sup>4</sup>The model in which the soft masses for the matter particles are generated from the nonvanishing part of the anomalous  $D$ -term was considered in [23] .

$$W_Y = 10 \cdot \bar{5} \cdot \bar{\phi} + \frac{Z_1}{M_P} 10 \cdot 10 \cdot \phi + \frac{Z_1}{M_P} \bar{5} \cdot 1 \cdot \phi , \quad (22)$$

where we have introduced  $Z$  and  $Z_1$  superfields with anomalous  $q_Z$  and  $q_{Z_1}$  charges respectively and assumed that  $Z$ ,  $Z_1$  share the VEVs with  $H$ ,  $\bar{H}$  and  $X$  fields in the  $D_A$  term. Therefore from (21) and (22)

$$\begin{aligned} q_{\bar{5}} &= -1 + \frac{1}{2}q_{Z_1} + q_Z , & q_{10} &= -1 - \frac{1}{2}q_{Z_1} , \\ q_1 &= -1 - \frac{3}{2}q_{Z_1} - q_Z . \end{aligned} \quad (23)$$

For  $1 - q_{Z_1}/2 < q_Z < -1 - 3q_{Z_1}/2$  and  $q_{Z_1} < -2$  charges of the all matter superfields will be positive.

Modification of the  $W_1^0$  do not change the picture of the SUSY breaking because the ratio of the ‘effective’  $F$  term and the mass of the messengers will be unchanged. As far as the masses of one pair of triplet-antitriplet components are concerned, their magnitudes are  $\sim M_X^2/M_P$ . This threshold will not spoil the picture of unification and even suggests the possibility of the obtaining small value of the  $\alpha_s$ . Namely, for  $\sin^2 \theta_W = 0.2313$ ,  $\alpha_s = 0.11$  is obtained <sup>5</sup>, which indeed coincides with the QCD sum analyses [29].

Until going to the fermion sector let us note that taking into account (15), (17) and (21) the nucleon decay parameter, which is the  $(1, 1)$  element of the inverse matrix of triplets-  $(\hat{M}_T^{-1})_{11}$  have the magnitude  $\sim m M_P / M_X^3$  and therefore nucleon decay through  $d = 5$  operator is strongly suppressed.

Turning to the fermion sector, we will see that  $R$  charges can be arranged in such a way that the  $\mu$ -term will have the desirable magnitude.

One of the nice feature of the flipped  $SU(5)$  theory is that in its framework, because  $e^c$  is identified as a singlet state of  $SU(5)$ , there do not exists the dangerous relation  $\hat{M}_d = \hat{M}_e$  which is concomitant to the minimal  $SU(5)$  theories. However, one can see that  $10 \cdot \bar{5} \cdot \bar{\phi}$  coupling generate the large “Dirac” mass for the neutrino. To suppress this mass the “Majorana” mass should be generated for the  $\nu_R$  state and then mass of the  $\nu_L$  can be suppressed by the universal seesaw mechanism [30].

Here we introduce the additional fermionic states  $\Psi \sim 10_1$ ,  $\bar{\Psi} \sim 10_{-1}$  and  $\mathcal{N} \sim 1_0$  (for each generation). Let us also introduce the three scalar superfields  $S$ ,  $S_1$  and  $S_2$ , which carry the  $U(1)_A$  charges. So they can also contribute in the  $D_A$ -term and can share the VEVs together with  $H$ ,  $\bar{H}$ ,  $X$ ,  $Z$  and  $Z_1$  states. So, if there do not exist for them couplings in the superpotential the abovepresented picture of the SUSY breaking will not changed. The transformation properties under the  $SU(5) \times U(1) \times U(1)_A \times \mathcal{R}$  symmetry of all introduced superfields are presented in the Table 1.

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<sup>5</sup>I thank I. Gogoladze for bringing my attention to this issue.



Table 1: The superfield transformation properties

Fields	$SU(5) \times U(1)$	$U(1)_A$	$\mathcal{R}$
$H$	$10_1$	$-1$	$R_{\bar{H}}$
$\bar{H}$	$\overline{10}_{-1}$	$-1$	$R_{\bar{H}}$
$\phi$	$5_{-2}$	$2$	$\alpha - 2R_H$
$\bar{\phi}$	$\bar{5}_2$	$2 - q_Z$	$\alpha - 2R_{\bar{H}} - R_Z$
$X$	$1_0$	$2$	$R_X$
$Y$	$1_0$	$0$	$R_Y$
$\bar{Y}$	$1_0$	$0$	$\alpha - R_H - R_{\bar{H}} - R_X - R_Y$
$Z$	$1_0$	$q_Z$	$R_Z$
$Z_1$	$1_0$	$q_{Z_1}$	$R_{Z_1}$
$S$	$1_0$	$q$	$R_S$
$S_1$	$1_0$	$q_1$	$R_1$
$S_2$	$1_0$	$2(q - q_1) - q_{Z_1}$	$-\alpha + 2(R_H + R_{\bar{H}} + R_S + R_X - R_1) - R_{Z_1}$
$10_i$	$10_1$	$-1 - \frac{1}{2}q_{Z_1}$	$R_H - \frac{1}{2}R_{Z_1}$
$\bar{5}_i$	$\bar{5}_{-3}$	$-1 + \frac{1}{2}q_{Z_1} + q_Z$	$-R_H + 2R_{\bar{H}} + R_Z + \frac{1}{2}R_{Z_1}$
$1_i$	$1_5$	$-1 - \frac{3}{2}q_{Z_1} - q_Z$	$3R_H - 2R_{\bar{H}} - R_Z - \frac{3}{2}R_{Z_1}$
$\Psi$	$10_1$	$1 - \frac{1}{2}q_{Z_1} + q - q_1$	$R_H + R_X + R_S - R_1 - \frac{1}{2}R_{Z_1}$
$\bar{\Psi}$	$10_{-1}$	$-1 + \frac{1}{2}q_{Z_1} - q$	$\alpha - R_H - R_X - R_S + \frac{1}{2}R_{Z_1}$
$\mathcal{N}$	$1_0$	$q_1 + \frac{1}{2}q_{Z_1} - q$	$\alpha - R_H - R_{\bar{H}} - R_X - R_S - R_1 + \frac{1}{2}R_{Z_1}$

Under these assignments of charges the Yukawa superpotential which generates masses of the quarks and leptons and also neutrinos "Dirac" and "Majorana" masses have the form:

$$\mathcal{W}_Y = W_Y + W'_Y , \quad (24)$$

where

$$W'_Y = A 10 \cdot \bar{\Psi} \cdot S \frac{X}{M_P} + B \bar{\Psi} \cdot \Psi S_1 + C \Psi \cdot \mathcal{N} \cdot \bar{H} + D S_2 \cdot \mathcal{N}^2 . \quad (25)$$

$A, \dots, D$  are the Yukawa matrices, which elements can be assumed to be of the order of one.

The neutrino mass matrix will have the form:

$$\hat{m}_\nu = \begin{matrix} & \nu_L & \nu_R & N_{\bar{\Psi}} & N_\Psi & N_{\mathcal{N}} \\ \begin{matrix} \nu_L \\ \nu_R \\ N_{\bar{\Psi}} \\ N_\Psi \\ N_{\mathcal{N}} \end{matrix} & \begin{pmatrix} 0 & h_u^0 & 0 & 0 & 0 \\ h_u^0 & 0 & \frac{XS}{M_P} & 0 & 0 \\ 0 & \frac{XS}{M_P} & 0 & S_1 & 0 \\ 0 & 0 & S_1 & 0 & \bar{H} \\ 0 & 0 & 0 & \bar{H} & S_2 \end{pmatrix} \end{matrix} . \quad (26)$$

This is the seesaw mass matrix which results the light neutrino with a mass of order:

$$m_\nu \simeq M_P^2 \frac{\langle h_u^0 \rangle^2}{M_X^3} . \quad (27)$$

Suppressing the neutrino masses the superpotential (25) was used in which the  $X$  superfield (with nonzero  $F$ -term) has couplings with superfields which are transforming nontrivially under the  $G_{321}$  gauge group and could emerge as a messengers.

Some remarks about implications of this fact should be done.

For simplicity let us consider the case of the one generation. After integrating out the heavy states of  $\bar{\Psi} + \Psi$ , which masses are  $\sim M_X$ , the decoupled state of decuplet (let us denote it by  $10_h$ ) will leave in 10 by the weight  $\frac{SX}{S_1 M_P} \sim \epsilon_X$ , while the light state (denoting it by  $10_f$ ) contained in it approximately by the weight 1. Taking into the account these facts the first two terms of the (25) can be rewritten as follows:

$$W_Y'^{(2)} = \frac{SX}{M_P} (\epsilon 10_h \cdot \bar{\Psi} + 10_f \cdot \bar{\Psi}) + S_1 \cdot \bar{\Psi} \cdot 10_h . \quad (28)$$

As we see the  $\bar{\Psi} + 10_h$  also emerge as a messengers but there contribution to the gaugino masses are strongly suppressed:

$$\delta M_a \sim \frac{\alpha_a}{4\pi} m \epsilon_X^2 . \quad (29)$$

Because there exists the matter coupling with messenger (see (28)) the soft masses for scalar components of  $10_f$  arise in the one loop level and as was shown in [31] is of the order:

$$\delta m_{10}^2 \sim \frac{1}{16\pi^2} \epsilon_X^2 \frac{F^4}{M^6} , \quad (30)$$

which is also negligible in comparison of (20).

Turning to the case of the three generations, without losing the generality the last term of (28) could be taken diagonal, while the remaining terms in general are nondiagonal. As was shown in [32] because of this fact the sparticles gain soft masses through the tadpole  $D$ -term of the  $U(1)_Y$ ; However in our case this contribution is miserable:

$$\delta m_{\tilde{\phi}_i}^2 \sim \frac{\alpha_1}{4\pi} Y_{\tilde{\phi}_i} m^2 \epsilon_X^4 . \quad (31)$$

We have not fixed yet the anomalous  $q$  and  $q_1$  as well as  $R_S$  and  $R_1$  charges of the  $S$  and  $S_1$  superfields respectively. Two cases can be considered:

**1.  $q = -5$  and  $R_S = -\frac{5}{2}R_X$**

In this case the term which is permitted by the  $U(1)_A$  and  $\mathcal{R}$  symmetries is

$$W_\mu = \bar{\phi}\phi \frac{X^3 Z Z_1 S_1^2 S_2}{M_P^7} . \quad (32)$$

Substituting VEVs  $\langle S_1 \rangle \sim \langle S_2 \rangle \sim \langle X \rangle \sim M_X$  and taking into account that  $M_X/M_P = \epsilon_X \sim 10^{-2}$  we will have  $\mu = M_X \epsilon_X^7 \sim 100$  GeV. Also substituting the nonzero  $F$ -term of the  $X$  superfield we will obtain the shift with the magnitude  $m M_X \epsilon_X^7$  between the masses of the scalar and fermionic components from  $\bar{\phi}$  and  $\phi$  superfields, this means that the standard doublets also emerge as messengers and wino can get mass through the one loop diagram [7]. In this case the masses of the gauginos will be:

$$M_3 = 4 \frac{\alpha_3}{4\pi} \frac{F_H}{\langle N \rangle} = \frac{\alpha_3}{\pi} m , \quad (33)$$

$$M_2 = 3 \frac{\alpha_2}{4\pi} \frac{m M_X \epsilon_X^7}{\mu} = 3 \frac{\alpha_2}{4\pi} m , \quad (34)$$

$$M_1 = \frac{\alpha_1}{4\pi} \left( 4 \frac{F_H}{\langle N \rangle} \frac{4}{30} + 3 \frac{m M_X \epsilon_X^7}{\mu} \frac{9}{30} \right) = \frac{43}{30} \frac{\alpha_1}{4\pi} m . \quad (35)$$

In (33) the factor 4 emerges because there are two pairs of the messengers and after the substitution of the  $F$ -terms in (5) the combinator factor 2 arise; while in (34) factor 3

arise because in (32) the field  $X$  (with nonzero  $F$ -term) is in the third power (the same arguments were taken into the account during the calculation of  $M_1$ ).

So, without introducing the additional states of the messenger superfields we can obtain the desirable pattern of SUSY breaking in the framework of the flipped  $SU(5)$  GUT, with successful DT splitting and with  $\mu$ -term of the order of 100 GeV. All messengers are from the minimal content of the GUT supermultiplets. Interestingly in this case the standard Higgs doublets also belong to the messenger superfields. Coupling of messenger doublets with ordinary matter induce the nonuniversalities in the one loop level but will be miserable in respect of (20) <sup>6</sup>.

**2.  $q = -\frac{4}{5}$  and  $R_S = -\frac{2}{5}R_X$**

In this case the  $\mu$ -term generating coupling is:

$$W_\mu = \bar{\phi}\phi \frac{S^3 Z Z_1 S_1^2 S_2}{M_P^7}, \quad (36)$$

which also gives  $\mu \sim 100$  GeV. In this case the doublets do not have couplings with superfields which  $F$ -terms have the nonzero VEVs.

Note, that in both cases the values of the  $q_1$  and  $q_Z$  charges still were undetermined. In order to insure the nonzero VEVs for  $H$ ,  $\bar{H}$  fields there charges and the  $\xi$  term in the  $D_A$  (see (8), (16)) must have the different signs. Therefore  $\text{Tr}Q = -46 - 28q_1 + 8q_Z - 21q_{Z_1}/2 > 0$  and from this condition we obtain  $q_1 < (16q_Z - 21q_{Z_1} - 92)/56$ .

Building our model we have assumed, that  $\epsilon_X = \sqrt{\xi}/M_P \sim 10^{-2}$  which for  $M_P \sim 10^{18}$  GeV gives  $\sqrt{\xi} \sim 10^{16}$  GeV, this value was dictated from the scale of the grand unification. However, for the flipped  $SU(5)$  model derived from strings [10] the preferable value of  $M_X$  is  $10^{17}$  GeV without losing the successful prediction of the  $\sin^2 \theta_W$  [33]. Increasing the value of  $M_X$  the picture of the abovepresented scenarios will not changed if for the values of  $q$ ,  $R_S$  for the two cases  $-13$ ,  $-13R_X/2$  and  $-4/13$ ,  $-2R_X/13$  will be taken respectively. Then the couplings generating the  $\mu$ -term will be  $X^{11} Z Z_1 S_1^2 S_2 / M_P^{15} \bar{\phi}\phi$  and  $S^{11} Z Z_1 S_1^2 S_2 / M_P^{15} \bar{\phi}\phi$  respectively, which for  $\epsilon_X = \sqrt{\xi}/M_P \sim 10^{-1}$  still give  $\mu \sim 100$  GeV.

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<sup>6</sup>See also the general arguments of ref. [9].

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